On the nucleon instability of heavy oxygen isotopes

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Abstract

The instability of the ^{26}O nucleus with respect to decay through the two-neutron channel is investigated. It is shown that this isotope, unobserved in the fragmentation experiments, can exist as a narrow resonance in the system $^{24}\mathrm{O} + 2n$. A role of deformation in formation of the neutron-drip line in the region $N \sim 20$ is discussed.

The structure of nucelei near the neutron-drip line possesses interesting features. These are: 1) two-neutron instability of nuclei which are stable with respect to one-neutron emission, 2) arising a new region of deformation when approaching the neutron-drip line, 3) change of the concept of nuclear shells at the boundary of neutron stability, 4) formation of the neutron halo, 5) clasterization in very neutron-rich nuclei. In this article, we pay attention to nucleon stability of oxygen nuclei near the neutron-drip line. This region, close to double magic area, has been attracting an attention of nuclear experimentalists an theorists for last decade.

The instability of the ^{26}O nucleus has been observed for the first time at the LISE spectrometer at GANIL in the fragmentation of the ^{48}Ca beam with an energy of 44 Mev per nucleon on a Ta target [1]. The work [2] confirmed this result. No events associated with the isotope ^{26}O were also observed in the recent analysis of oxygen fragments by the separator RIPS in the projectile fragmentation experiments with using 94.1 MeV per nucleon ^{40}Ar beam at the RIKEN [3]. Theoretical study of properties of ^{26}O was done in [4] within the framework of the self-consistent theory of finite Fermi systems [5, 6]. Estimation of one-neutron separation energy yielded $S_n \simeq 0.7\,\mathrm{MeV}$, i.e. this nucleus was found to be stable to one-neutron emission. This clears the way for the pure two-neutron instability of the ^{26}O nucleus. Before going over to investigation of this subject we concern briefly neighbouring even-even oxygen isotopes. Calculations made in [4] for the nucleus ^{24}O observed in all the experiments yielded $S_n = 3.6\,\mathrm{MeV}$ and $S_{2n} = 5.9\,\mathrm{MeV}$. At the same time those calculations showed $1n$- and $2n$-instability of the nucleus ^{28}O unobserved in [1, 2, 3].

The two-neutron instability corresponding to positive one-neutron separation energy S_n and negative two-neutron separation energy S_{2n} is possible owing to the

positive energy of neutron pairing [9]. This energy is usually characterized by the even-odd difference Δ_n of the binding energies E_b of three neighboring isotopes, which is given by

$$\Delta_n = \frac{1}{2}\Big[E_b(N{+}2,Z) - 2E_b(N{+}1,Z) + E_b(N,Z)\Big] \tag{1}$$

and is positive for even-even nucleus (N,Z). The obvious relation $S_{2n}-2S_n = -2\Delta_n$ follows from eq. (1). This relation indicates that two-neutron instability is possible provided that the condition $0{<}S_n{<}\Delta_n$ is satisfied. Bearing in mind that in this region of nuclei $\Delta_n \sim 2 \div 3\,\mathrm{MeV}$ one can see that this condition holds for the ^{26}O nucleus.

To investigate the stability of the ^{26}O nucleus, we consider the $^{24}\mathrm{O}+2n$ system. The energies E_q and widths Γ_q of quasistationary states of this system are determined by the poles $E_q - i\Gamma/2$ of the Green's function of two neutrons in the mean field of the ^{24}O nucleus [5]. The equation for the two-neutron Green's function can be reduced to the equation for the two-neutron wave function $\Psi(1,2)$:

$$\Big[H_0(1) + H_0(2) + V(1,2)\Big]\Psi(1,2) = E\Psi(1,2), \tag{2}$$

where the arguments 1 and 2 stand for the sets of neutron spatial and spin coordinates, H_0 is the single-particle Hamiltonian, and V is the effective interaction in the two-particle channel. Recent microscopic calculation of this interaction for semi-infinite nuclear matter [7] showed its surface nature. The phenomenological analysis of even-odd effects in tin isotopic chain [8] showed that density dependence of the effective pairing interaction is quite sophisticated. However the parameters controlling the shape and the strength of this interaction are fitted for rather heavy nuclei not very far from the valley of β-stability and can hardly be directly used in the problem under discussion. That is why we did not claim to calculate the wave function $\Psi(1,2)$ and the energy E. The aim is to estimate the width of the two-neutron state by the order of value.

To solve eq. (2), we expand the function $\Psi(1,2)$ in the basis formed by the products of two eigenfunctions of the Hamiltonian H_0 with the total angular momentum $I{=}0$ and its projection $M{=}0$. This expansion has the form

$$\Psi(1,2) = \int_0^\infty\!\!\int_0^\infty d\varepsilon_1 d\varepsilon_2 C^\nu(\varepsilon_1,\varepsilon_2)\Big\{\psi^\nu_{\varepsilon_1}(1)\cdot\psi^\nu_{\varepsilon_2}(2)\Big\}^{00}, \tag{3}$$

where ν is the set of angular quantum numbers.

For low energies, the ε-dependence of the functions $\psi_\varepsilon(1)$ can be separated from the coordinate and spin dependences as follows [10, 11]:

$$\psi_\varepsilon(1) = \sqrt{\Delta(\varepsilon)}\psi_0(1). \tag{4}$$

According to the theoretical scheme of neutron single-partcle levels of the ^{24}O nucleus, the $2s_{1/2}$ state is the shallowest bound state, and the $1d_{3/2}$ state is the lowest quasistationary state. At low energies, the factor

$$\Delta^d(\varepsilon) = \frac{1}{2\pi}\,\frac{\gamma}{(\varepsilon-\varepsilon_0^d)^2+\gamma^2/4} \tag{5}$$

is resonantly enhanced for the $d_{3/2}$ continuum states because of proximity to the $1d_{3/2}$ quasistationary state with energy $\varepsilon^d=\varepsilon_0^d{-}i\gamma/2$ in the ^{24}O nucleus [10]. Owing to existence of the weakly bound $2s_{1/2}$ state, the wave functions of the $s_{1/2}$ continuum states are enhanced at low energies by the factor [11]

$$\Delta^s(\varepsilon) \simeq \frac{\sqrt{2m}\beta}{\pi\sqrt{\varepsilon}(\varepsilon{+}\beta)}, \tag{6}$$

where $\beta{=}2/mR^2$ and R is of the order of the ^{24}O radius. For this reasons, we restrict the basis in the expansion (3) of the wave function $\Psi(1,2)$ to the $s_{1/2}$ and $d_{3/2}$ states of single-particle continuum.

To derive the required relations, we follow the method used in [10] to obtain the relations of the theory of two-proton radioactivity. Substituting expansion (3) with the functions $\psi_\varepsilon(1)$ factorized in form (4) into eq. (2), we find that the energy of the quasistationary state satisfies the algebraic equation

$$g_{dd}\,M^d(E)+g_{ss}\,M^s(E)+(g_{ds}^2{-}g_{dd}\,g_{ss})\,M^d(E)\,M^s(E)=1, \tag{7}$$

where

$$M^\lambda(E)=\int\limits_0^\infty\!\int\limits_0^\infty d\varepsilon_1 d\varepsilon_2\,\frac{\Delta^\lambda(\varepsilon_1)\Delta^\lambda(\varepsilon_2)}{\varepsilon_1{+}\varepsilon_2{-}E}, \tag{8}$$

$$g_{\lambda\nu}=-\int d1\,d2\,V(1,2)\left\{\psi_0^\lambda(1)\cdot\psi_0^\lambda(2)\right\}^{00}\left\{\psi_0^\nu(1)\cdot\psi_0^\nu(2)\right\}^{00}, \tag{9}$$

and $\lambda,\nu=s,d$. Substituting expresiions (5) and (6) for the quantity $\Delta^\lambda(\varepsilon)$ into eq. (8) and calculating the integrals by rotating the contour to the negative imaginary axis, we obtain

$$M^s(E)=2m\left(-\frac{1}{\pi}\ln\frac{E}{4\beta}-\frac{1}{2}+i\right), \tag{10}$$

$$M^d(E)=\frac{1}{2\varepsilon_0^d{-}E}+\frac{2\gamma^2}{\pi(2\varepsilon_0^d{-}E)^2}\left\{\frac{E}{\varepsilon_0^d{-}E}+\frac{1}{2\varepsilon_0^d{-}E}\ln\left|\frac{\varepsilon_0^d}{\varepsilon_0^d{-}E}\right|\right\}. \tag{11}$$

Eq. (7) with $M^s(E)$ and $M^d(E)$ specified by eqs. (10) and (11) is an algebraic equation with complex coefficients. Numerical analysis of eq. (7) shows that it has

a complex root $E = E_0 - i\Gamma/2$ in a wide region of input parameters. While the position of this root depends on the values of the parameters, it is located around $E_0 \sim 0.1\,\mathrm{MeV}$, $\Gamma \sim 10^{-3}\,\mathrm{MeV}$. We reserve a detailed study of this dependence for the future publication. The width $\Gamma \sim 10^{-3}\,\mathrm{MeV}$ of this quasistationary state corresponds to a lifetime $\tau \sim \hbar/\Gamma \sim 10^{-18}\,\mathrm{s}$. Since $\tau \gg \tau_{nucl} \sim 10^{-22}\,\mathrm{s}$ the emitted neutron pair should be a weakly bound di-neutron state. It can be observed in correlation experiments. The similar situation was discussed earlier in connection with an analyzis of the β-delayed multi-neutron emission [12, 13]. However, in that case the cascade $(n+n)$ neutron emission strongly dominates and the di-neutron (^2n) channel is suppressed (e.g. for ^{35}Na, $P_{^2n}/P_{n+n} < 0.19$ [13]) due to a decay of the excited state. In the case of ^{26}O, the neutron pair is emitted from the ground state, so that the situation is "more clear".

The spherical basis was used in calculations for neutron-rich oxygen isotopes. However, account of deformation (even small, with $\beta_2 < 0.2$) could strongly change the picture. Indeed, splitting of the single-particle $d_{3/2}$ neutron state results in destroying the shell $N = 20$ for nuclei near the neutron-drip line [15]. Such a sensitivity to small deformations should essentially complicate description of neutron-rich nuclei. However, the oxygen isotopes evidently should not be deformed. Otherwise, being deformed with $\beta_2 > 0.1$, the isotope ^{26}O could obtain additional stability due to increase of the binding energy by ~ 1 MeV and should be observed experimentally, but this is not the case.

This is the deformation that seems to give an explanation of existence of the isotope ^{31}F observed in the experiment [3]. The calculations [4] of one- and two-neutron separation energies for this nucleus based on the self-consistent finite Fermi system theory in the sphericl geometry yielded nucleon instability of ^{31}F. At the same time the scenario of sharp shift of the neutron-drip line owing to onset of a deformation was suggested several years ago for explanation of the nucleon stability of heavy sodium isotopes [14]. Nuclei with less than half occupation of the level $f_{7/2}$ being slightly unbound in spherical calculation ($S_n \lesssim 0$) can become stable due to lowering the energy of states with asymptotic quantum numbers $\frac{1}{2}^-[330]$ and $\frac{3}{2}^-[321]$ at deformation [16]. The analogous scenario seems to take place for fluorine isotopes: the neutron level $f_{7/2}$ just starts to be occupied in the nucleus ^{31}F with $N = 22$. Following this scenario one could expect nucleon stability of the isotope ^{33}F as well. The problem of onset of a deformation for nuclei near the neutron-drip line will be discussed in a separate article, in particular, in connection with an opportunity of clasterization in weakly bound neutron-rich systems.

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